

Flavored Gauge Mediation, A Heavy Higgs, and Supersymmetric Alignment

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Abstract

We show that the messenger-matter couplings of Flavored Gauge Mediation Models can generate substantial stop mixing, leading to Higgs masses around 126 GeV with colored superpartners below 2 TeV and even a TeV. These results are largely independent of the messenger scale. We study the spectra of a few examples with a single messenger pair coupling dominantly to the top, for different messenger scales. Flavor constraints in these models are obeyed by virtue of supersymmetric alignment: the same flavor symmetry that explains fermion masses dictates the structure of the matter-messenger couplings, and this structure is inherited by the soft terms. We also present the leading 1-loop and 2-loop contributions to the soft terms for general coupling matrices in generation space.

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I. INTRODUCTION

The null results of direct searches for supersymmetry suggest that it does not manifest itself in the form of light, flavor-blind superpartners. Furthermore, if the recently discovered scalar at 126 GeV [1, 2] is indeed the Higgs boson, its relatively large mass requires either a large stop mixing or very heavy stops, at least in the context of the Minimal Supersymmetric Standard Model (MSSM) [3]. These results are especially problematic for Gauge-Mediated Supersymmetry Breaking (GMSB) models [4, 5], which predict zero A -terms and flavor-blind spectra at the messenger scale. Low scale gauge mediation is therefore strongly disfavored by the Higgs mass, and even high-scale models, with A -terms generated by the running below the messenger scale, require stop masses of around 8–10 TeV [6, 7] in the context of minimal gauge mediation [5]. Given the tight relations between squark and gluino masses in minimal GMSB (mGMSB), this implies that all colored superpartners are very heavy in these models, and beyond the reach of any foreseeable experiment.

From a purely theoretical point of view, however, GMSB models are very attractive, since both the breaking of supersymmetry and its mediation are described by well understood quantum field theories, as opposed to unknown Planck-scale physics. Indeed, flavor-blind extensions of gauge mediation have been extensively discussed in recent years [8]. These extensions too are only consistent with a 126 GeV Higgs for high messenger scales, unless the stops or the gluino are very heavy [9]. Here we will study instead a flavor-dependent extension of gauge mediation, specifically, the Flavored Gauge Mediation (FGM) models of [10]. In these models, the flavor structure of GMSB is in principle modified, due to superpotential couplings of the messengers to SM fields [11]. We will show that these messenger-matter couplings can yield significant top A -terms, and therefore a heavy Higgs, even for fairly light superpartners, for a wide range of messenger scales.

The superpartner masses in FGM models are generated by both the SM gauge interactions, and by Yukawa-type superpotential couplings of the messengers to the SM matter fields. Thus, while the interactions mediating the breaking are not purely gauge interactions, they are still completely “visible”—occurring within simple field theoretic extensions of the MSSM, and potentially at low scales. Since the matter-messenger couplings are in principle flavor dependent, they are strongly constrained by the non-observation of flavor changing neutral currents. As stressed in [10], however, there are good reasons to consider these couplings. At the very least, given our ignorance about the origin of the SM Yukawas structure, it is conceivable that the messenger-matter couplings have some special structure, which results in an acceptable pattern of soft terms. Indeed, the structure of the SM fermion masses hints at some theory of flavor, and any such theory will necessarily also control the sizes of matter-messenger couplings. Furthermore, as superpartner masses are pushed to higher values by direct searches (as well as by the large Higgs mass), there is more room for non-degenerate spectra. From the point of view of LHC searches for supersymmetry, the assumption of Minimal Flavor Violation (MFV), which underlies many analyses, can result in reduced sensitivity to non-MFV spectra [12–14]. So when searching for gauge-mediated supersymmetry, it is important to keep the possibility of flavor-dependent spectra in mind, and FGM models provide useful examples of such spectra.

The main new ingredient in our models will be superpotential messenger-matter interactions, with the up-type Higgs (or also the down-type Higgs) replaced by a messenger field of the same gauge charges. Since we would like to generate a large top A -term, the messengers need only couple to the top. As a concrete realization of this scenario, we in-

voke a flavor symmetry under which the Higgs and messenger field have identical charges. Flavor constraints in our models are thus satisfied by a combination of degeneracy, coming from the pure gauge contributions, and alignment [10]. Unlike in the original alignment models of [15], in which the flavor symmetry controls the soft terms directly, here it controls the *supersymmetric* messenger-matter couplings so that they are aligned with the SM Yukawas. The soft terms therefore inherit this structure, even though they are generated at the messenger scale, which is typically much lower than the flavor-symmetry breaking scale.

We note that three other papers appeared recently [16–18] which rely on messenger-SM couplings to raise the Higgs mass. The differences between our models and the models of [16–18] arise due to the different choices of symmetries and, as a result, the allowed messenger-SM couplings. In [16], the messengers are chosen to have the same R-parity as the SM matter fields, so the relevant coupling is the analog of the Yukawa coupling with one SM matter field replaced by a messenger. Messengers in $5+\bar{5}$ representations of $SU(5)$ in these models do not affect the u^c mass, and as a result can raise the Higgs mass only to around 118 GeV for stops below 1.5 TeV. Therefore [16] uses messengers in $10+\bar{10}$. In [17, 18], a coupling of the type Higgs-messenger-messenger is used, with one messenger being a SM gauge singlet. Since none of the fields involved is colored, the effect of this coupling is moderate, so that the Higgs mass is viable only at low messenger scales, where the negative one-loop contributions to the Higgs soft masses-squared are important [18]. It is interesting that, although the new messenger couplings in our models do not involve the Higgs at all, they have a significant effect on the Higgs mass. The reason is that the key feature needed for getting a large Higgs mass without resorting to large stop masses, is the top A -term. This only requires that the messengers couple to the top, and since the new couplings then involve colored particles, the effect on the Higgs mass can be dramatic.

The organization of this paper is as follows. In Section II A we briefly review FGM and introduce the symmetries and the superpotential of our models. In Section II B we give expressions for the soft terms in the limit of third-generation dominance¹. In Section III we present the Higgs mass and superpartner spectra for different choices of parameters. Our conclusions and discussion of the results are presented in Section IV.

II. MODELS

A. The models and supersymmetric alignment.

We begin by briefly reviewing FGM models [10]. The starting point in these models is mGMSB [5]. Specifically, we will take N sets of messengers transforming as $5+\bar{5}$ of $SU(5)$, coupled to a supersymmetry-breaking singlet $\langle X \rangle = M + F\theta^2$. We use capital letters to denote the messenger fields, with $5 = T + D$ and $\bar{5} = \bar{T} + \bar{D}$, where T (\bar{T}) and D (\bar{D}) are fundamentals (anti-fundamentals) of $SU(3)$ and $SU(2)$ respectively. The SM gauge symmetry permits different couplings of the messengers to SM fields, and these would generically give rise to flavor-dependent soft-terms [11, 19–22]².

As in [10], we will assume that the SM fermion masses are explained by a flavor symmetry.

¹ Full expressions for the soft terms for general 3×3 coupling matrices in generation space are presented in the Appendix.

² The model of [20] relies on an extra dimension in order to obtain MFV couplings.

This symmetry then also controls the messenger-matter couplings³. In the models we will consider, these coupling matrices will be aligned with the SM Yukawa matrices, so that flavor constraints are satisfied. Naively, one would think that alignment can only be relevant for high-scale supersymmetry breaking. This is indeed the case in the original alignment models of [15]. In these models, a Froggatt-Nielsen flavor symmetry [23] dictates the structure of the soft-term matrices at the supersymmetry-breaking scale. As explained above however, the non-universal parts of the soft terms in FGM models are generated by superpotential matter-messenger couplings. These *supersymmetric* coupling matrices are the ones controlled by the flavor symmetry at high scales, and their near-diagonal structure is inherited by the soft terms, which are generated at much lower scales. We therefore refer to this type of alignment as “supersymmetric alignment”.

In addition to the flavor symmetry and R-parity, we impose a Z_3 symmetry with charges given in Table I. The following superpotential is then allowed by the symmetries,

Superfield	R -parity	Z_3
X	even	1
D_1	even	-1
\bar{D}_1	even	0
D_2	even	0
\bar{D}_2	even	-1
$T_I, \bar{T}_I, D_{I>2}, \bar{D}_{I>2}$	even	1
q, u^c, d^c, l, e^c	odd	0
H_U, H_D	even	0

TABLE I: R-parity and Z_3 charges.

$$W = X (X^2 + T_I \bar{T}_I + D_I \bar{D}_I) + Y_U H_U q u^c + Y_D H_D q d^c + Y_L H_D l e^c + y_U \bar{D}_1 q u^c + y_D D_2 q d^c + y_L D_2 l e^c, \quad (1)$$

where $I = 1, \dots, N$ runs over messenger pairs, y_U, y_D, y_L are messenger-matter Yukawa matrices, Y_U, Y_D, Y_L are the SM Yukawa matrices, and q, u^c, d^c, l, e^c are the MSSM chiral multiplets. We assume that the μ -term(s) are forbidden by some U(1) symmetry. Note that to have messenger couplings to both up quarks and down quarks we need at least two sets of messengers [10, 20]. We also display here the term X^3 , required in mGMSB in order to generate the X VEVs, and motivating our choice of a Z_3 symmetry. In the following, however, we will limit ourselves to treating X as a supersymmetry breaking spurion.

At this point, D_2 and H_D , as well as \bar{D}_1 and H_U , have identical charges under all the symmetries, and therefore the following terms are allowed as well,

$$X D_1 H_U + X H_D \bar{D}_2. \quad (2)$$

However, we can set these couplings to zero without loss of generality. Consider for concreteness the H_U and \bar{D}_1 couplings

$$y_{Uij} \bar{D}_1 q_i u_j^c + Y_{Uij} H_U q_i u_j^c, \quad (3)$$

³ In general, some messenger fields may be charged under the flavor symmetry.

where i, j are generation indices. Taking H_U and \bar{D}_1 to have the same charges under the flavor symmetry, we can define the combination of \bar{D}_1 and H_U that couples to X to be the messenger (indeed, this is the massive eigenstate), and the orthogonal combination to be the Higgs. A similar redefinition can be done for D_2 and H_d . Thus (1) is the most general superpotential and the entries y_{Uij} and Y_{Uij} are the same up to order-one coefficients⁴. Since the only order-one entry of Y_U is Y_{U33} , the two matrices Y_U and y_U are approximately diagonal, and the soft terms inherit this structure. Inter-generational mixings are thus suppressed by supersymmetric alignment⁵.

B. A -terms and scalar masses

At leading order in F/M^2 , the messenger-matter couplings of Eq. (1) generate one-loop contributions to the A -terms, and two-loop contributions to the sfermion masses-squared. We present full expressions for the soft terms in Appendix A. In the case of interest, only the 3-3 entries of the coupling matrices are important and the soft terms (at the messenger scale) simplify to,

$$\begin{aligned} A_{33}^U &= -\frac{Y_t}{16\pi^2} [3y_t^2 + y_b^2] \frac{F}{M} \\ A_{33}^D &= -\frac{Y_b}{16\pi^2} [3y_b^2 + y_t^2] \frac{F}{M} \\ A_{33}^L &= -\frac{3Y_\tau y_\tau^2}{16\pi^2} \frac{F}{M} \end{aligned} \tag{4}$$

⁴ The running between the UV scale and the messenger scale will introduce, of course, some mixing between H_U and \bar{D}_1 , but the only effect of this running is to modify the order-one coefficients of y_U and Y_U .

⁵ It is also possible to choose different charges for H_U and \bar{D}_1 , (and similarly for H_D and D_2) such that the terms (2) are either forbidden or very suppressed. In this case, y_U and Y_U will have different textures, and these can be chosen to be compatible with flavor constraints [10].

where $Y_t \equiv Y_{U_{33}}$ and similarly for the remaining couplings, and,

$$\begin{aligned}
\tilde{m}_{H_U}^2 &= \frac{1}{128\pi^4} \left\{ -\frac{3}{2}Y_t^2(3y_t^2 + y_b^2) + N \left(\frac{3}{4}g_2^4 + \frac{3}{20}g_1^4 \right) \right\} \left| \frac{F}{M} \right|^2 \\
\tilde{m}_{H_D}^2 &= \frac{1}{128\pi^4} \left\{ -\frac{3}{2}Y_b^2(3y_b^2 + y_t^2) - \frac{3}{2}Y_\tau^2 y_\tau^2 + N \left(\frac{3}{4}g_2^4 + \frac{3}{20}g_1^4 \right) \right\} \left| \frac{F}{M} \right|^2 \\
(\tilde{m}_q^2)_{33} &= \frac{1}{128\pi^4} \left\{ \left(y_t^2 + 3y_b^2 + \frac{1}{2}y_\tau^2 - \frac{8}{3}g_3^2 - \frac{3}{2}g_2^2 - \frac{7}{30}g_1^2 \right) y_b^2 \right. \\
&\quad \left. + \left(3y_t^2 - \frac{8}{3}g_3^2 - \frac{3}{2}g_2^2 - \frac{13}{30}g_1^2 \right) y_t^2 + N \left(\frac{4}{3}g_3^4 + \frac{3}{4}g_2^4 + \frac{1}{60}g_1^4 \right) \right\} \left| \frac{F}{M} \right|^2 \\
(\tilde{m}_{u^c}^2)_{33} &= \frac{1}{128\pi^4} \left\{ \left(6y_t^2 + y_b^2 + Y_b^2 - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{13}{15}g_1^2 \right) y_t^2 - Y_t^2 y_b^2 \right. \\
&\quad \left. + N \left(\frac{4}{3}g_3^4 + \frac{4}{15}g_1^4 \right) \right\} \left| \frac{F}{M} \right|^2 \\
(\tilde{m}_{d^c}^2)_{33} &= \frac{1}{128\pi^4} \left\{ \left(6y_b^2 + y_\tau^2 + y_t^2 + Y_t^2 - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{7}{15}g_1^2 \right) y_b^2 - y_t^2 Y_b^2 \right. \\
&\quad \left. + N \left(\frac{4}{3}g_3^4 + \frac{1}{15}g_1^4 \right) \right\} \left| \frac{F}{M} \right|^2, \\
(\tilde{m}_l^2)_{33} &= \frac{1}{128\pi^4} \left\{ \left(\frac{3}{2}y_b^2 + 2y_\tau^2 - \frac{3}{2}g_2^2 - \frac{9}{10}g_1^2 \right) y_\tau^2 + N \left(\frac{3}{4}g_2^4 + \frac{3}{20}g_1^4 \right) \right\} \left| \frac{F}{M} \right|^2 \\
(\tilde{m}_{e^c}^2)_{33} &= \frac{1}{128\pi^4} \left\{ \left(3y_b^2 + 4y_\tau^2 - 3g_2^2 - \frac{9}{5}g_1^2 \right) y_\tau^2 + \frac{3}{5}N g_1^4 \right\} \left| \frac{F}{M} \right|^2.
\end{aligned} \tag{5}$$

If the messenger scale M is below roughly 10^7 GeV, the one-loop $\mathcal{O}(F^4/M^6)$ contributions [11] to the soft masses may be important. In the limit of third-generation dominance, these contributions are given by

$$\begin{aligned}
(\delta\tilde{m}_q^2)_{33} &= -\frac{1}{6} \frac{1}{16\pi^2} (y_t^2 + y_b^2) \frac{F^4}{M^6} \\
(\delta\tilde{m}_{u^c}^2)_{33} &= -\frac{1}{3} \frac{1}{16\pi^2} y_t^2 \frac{F^4}{M^6} \\
(\delta\tilde{m}_{d^c}^2)_{33} &= -\frac{1}{3} \frac{1}{16\pi^2} y_b^2 \frac{F^4}{M^6} \\
(\delta\tilde{m}_l^2)_{33} &= -\frac{1}{6} \frac{1}{16\pi^2} y_\tau^2 \frac{F^4}{M^6} \\
(\delta\tilde{m}_{e^c}^2)_{33} &= -\frac{1}{3} \frac{1}{16\pi^2} y_\tau^2 \frac{F^4}{M^6},
\end{aligned} \tag{6}$$

While the next-to-leading contribution in F/M^2 to the top A -term is

$$\delta A_{33}^U = -\frac{Y_t y_t^2}{16\pi^2} \frac{F^3}{M^5}. \tag{7}$$

Comparing the new contributions to the mGMSB expressions, we see that the importance of the new contributions relative to the mGMSB expressions is maximal for the smallest

possible number of messengers. Thus we will take $N = 1$ when the only new Yukawa coupling is y_t , and $N = 2$ if y_b and/or y_τ are present as well.

The messenger-scale A -terms are the crucial new ingredient relevant for the Higgs mass. As is well known [3], at one-loop, the Higgs mass is approximately given by,

$$m_h^2 = m_Z^2 \cos^2 2\beta + \frac{3m_t^4}{4\pi^2 v^2} \left[\log \frac{M_S^2}{m_t^2} + \frac{X_t^2}{M_S^2} \left[1 - \frac{X_t^2}{12M_S^2} \right] \right], \quad (8)$$

where $X_t = A_t/Y_t - \mu \cot \beta$ is the LR stop mixing and $M_S \equiv (m_{\tilde{t}_1} m_{\tilde{t}_2})^{1/2}$ is the average stop mass. Clearly, the Higgs mass can be increased either by increasing the average stop mass, so that the log term is large, or by increasing the stop mixing, so that X_t/M_S is large, with the maximal m_h^2 obtained for X_t/M_S of around 2.4 [24].

Since our main objective is to obtain the correct Higgs mass with superpartners within the LHC reach, we need a large X_t/M_S , and therefore a large A_t . As can be seen from equations (5), a non-zero y_t , which generates A_t , also gives new contributions to the stop masses. This latter contribution is positive for y_t much larger than g_3 , however, it is negative for intermediate values of y_t . Thus one may obtain significant values of X_t/M_S for some range of coupling y_t . It is natural to expect this region to occur for y_t near g_3 , which is order one, and as we will see in explicit examples below, this is indeed the case. This is very fortunate, because, as discussed above, the flavor symmetry we chose implies values of y_t that are of order one.

III. HIGGS MASS AND SUPERPARTNER MASSES

A. The Higgs and stop masses

We first consider models with one set of messengers. With the Z_3 charges of Table I, only the \bar{D}_1 messenger couplings to matter are allowed. Moreover, since we assume that \bar{D}_1 and H_U have the same charge under the flavor symmetry of the model, the only significant messenger coupling is $y_t \equiv (y_U)_{33}$. We use SOFTSUSY [25] to calculate the Higgs mass for different choices of the GMSB parameters and y_t . Given the theoretical uncertainty in the calculation of the Higgs mass, it is interesting to study Higgs masses in the 124–128 GeV window.

In Fig. 1 we show contours of the Higgs and stop masses as a function of $\Lambda \equiv F/M$ and y_t , for a low messenger scale of $M = 900$ TeV, with $\tan \beta = 10$. For such a low messenger scale, the one-loop $\mathcal{O}(F^4/M^6)$ corrections are not necessarily negligible and have been taken into account. Fig. 1(a) shows the Higgs mass contours for a wide range of y_t . This allows us to compare the Higgs mass prediction in FGM models to minimal gauge mediation: the mGMSB predictions can be read off at $y_t = 0$. The white region for intermediate values of y_t is excluded. In this region, the stops are either tachyonic or too light for successful electroweak symmetry breaking. As explained above, for these values of y_t , the negative $y_t^2 g^2$ contribution to the stop mass-squared is comparable to the positive contributions from pure GMSB. As y_t is increased further, the y_t^4 contribution to the stop masses guarantees that the stops are non-tachyonic, but because of the partial cancellation between the $y_t^2 g^2$ term and the remaining terms, the A -term becomes appreciable compared to the stop masses and the resulting large mixing allows for a heavy Higgs. As y_t is increased even further, the

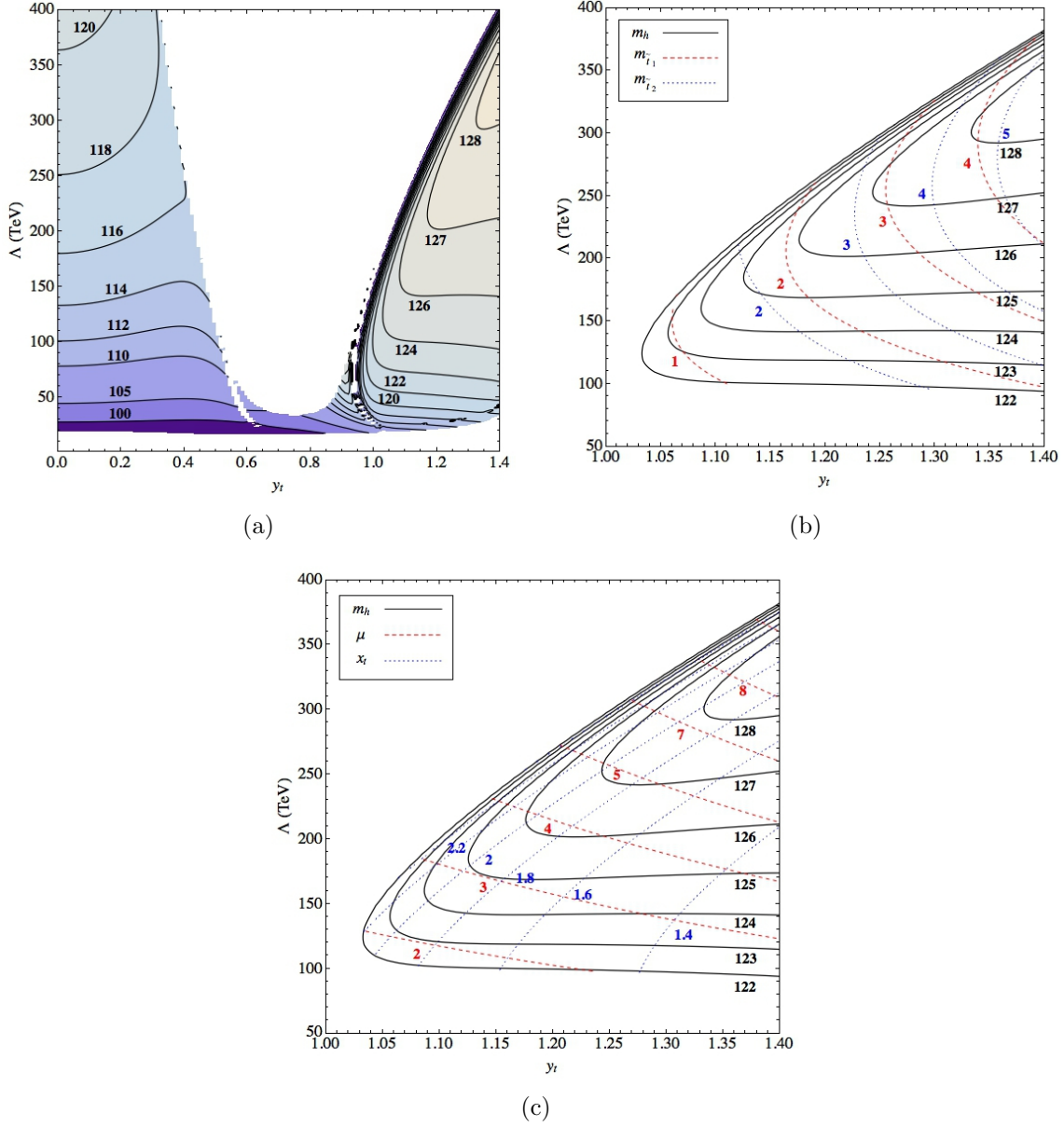


FIG. 1: The Higgs and stop masses for $N = 1$, $M = 900$ TeV, $\tan \beta = 10$. Fig. 1(a) shows the Higgs mass for a wide range of y_t . The predictions of minimal gauge mediation can be read off from the line $y_t = 0$. The white region is excluded because it leads to tachyonic stops (see text). In Fig. 1(b), we show Higgs mass (solid), heavy stop mass (dotted) and light stop mass (dashed) contour lines in a smaller region y_t . In Fig. 1(c) we show Higgs (solid), μ (dashed) and $x_t = |X_t/M_S|$ (dotted) contour lines in the same region.

stops become heavy relative to the other superpartners. However, even in this regime the A-terms play a crucial role in making the Higgs heavy ⁶.

⁶ To understand the behavior of the Higgs mass for small values of y_t , note that in this regime the Higgs mass is mainly affected by the correction to the stop mass, which is dominated by the negative $y_t^2 g^2$ term.

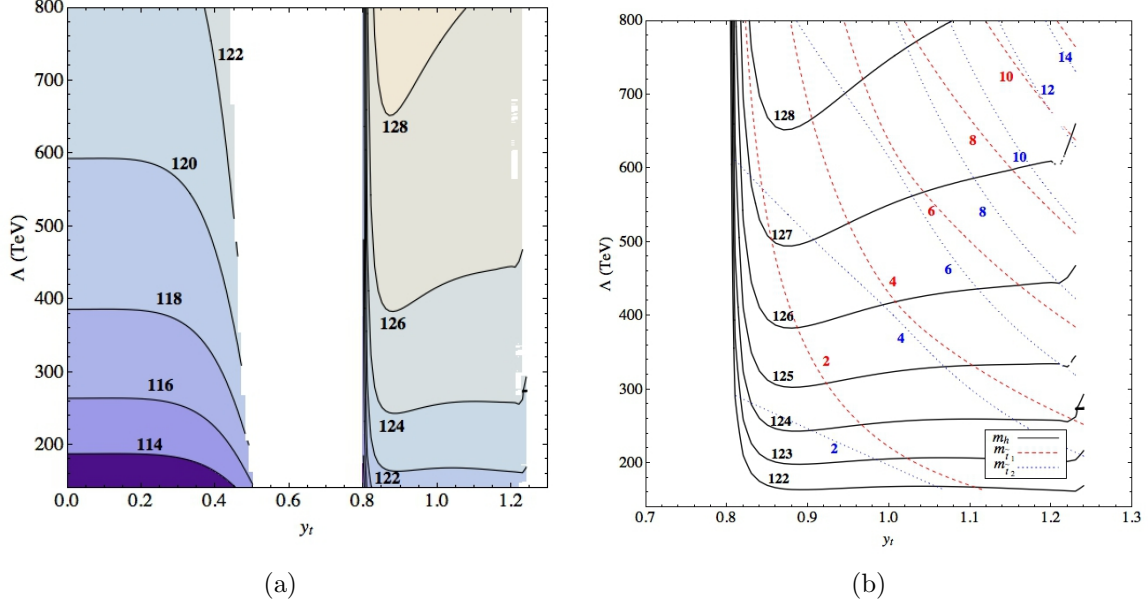


FIG. 2: Same plots as in Fig 1(a) and 1(b) with $M = 10^{12}$ GeV and $\tan \beta = 10$

In Fig. 1(b), we zoom in on the interesting range $y_t \sim 1$, and show contours of the Higgs mass together with the two stop masses, with the remaining parameters being the same as in Fig. 1(a). In Fig. 1(c), we also show contours of μ and the mixing parameter $x_t = |X_t/M_S|$. As expected, the largest values of x_t are obtained close to the excluded white region, where there is a partial cancellation between the different contributions to the stop masses. Thus, appreciable A -terms can be obtained without a large increase in the stop squared masses. Indeed, as can be seen in Fig. 1(b), for these large values of x_t , the Higgs mass can be large even for low Λ 's, such that the stops are light. For values of y_t of order one we can therefore find at least one stop below 2 TeV, and, when the accidental cancellation between the g^4 and $g^2 y_t^2$ terms is significant, even below a TeV.

The qualitative behavior of the results is not very sensitive to the value of the messenger scale. To demonstrate this, we present similar plots for two other messenger scales, $M = 10^{12}$ GeV in Fig 2 and $M = 10^8$ GeV in Fig 3. In both cases a moderate messenger Yukawa, $y_t \sim 0.8 - 0.9$, may be sufficient to obtain Higgs masses above 123 GeV. It is especially useful to compare our results for high messenger scales with models of minimal gauge mediation. As is well known, with a lot of running, appreciable A -terms can be generated in pure GMSB models. This is, however, not sufficient — as was shown in Ref. [7], even with a high messenger scale a heavy Higgs requires very heavy stops near 8-10 TeV. For example, with $M = 7.9 \cdot 10^{12}$ GeV and $\tan \beta = 10$, a Higgs mass of 125 GeV can be achieved if $\Lambda = 1.3 \cdot 10^6$ GeV [7]. With such a high value of Λ , one of the stops is the lightest squark and has a mass of 7.9 TeV. In contrast, if we choose the same messenger scale in the FGM model with $y_t = 0.78$, a Higgs mass of 127 GeV can be achieved with $\Lambda = 2.4 \cdot 10^5$ GeV. This value of Λ results in the first two generation squarks being around 2 TeV. The stop masses are lowered even further due to the negative $y_t^2 g^2$ contributions, with the lightest stop mass

For fixed Λ , this negative contribution would result in a smaller M_S . Therefore, to achieve the same Higgs mass, larger values of Λ are required compared to the case of mGMSB.

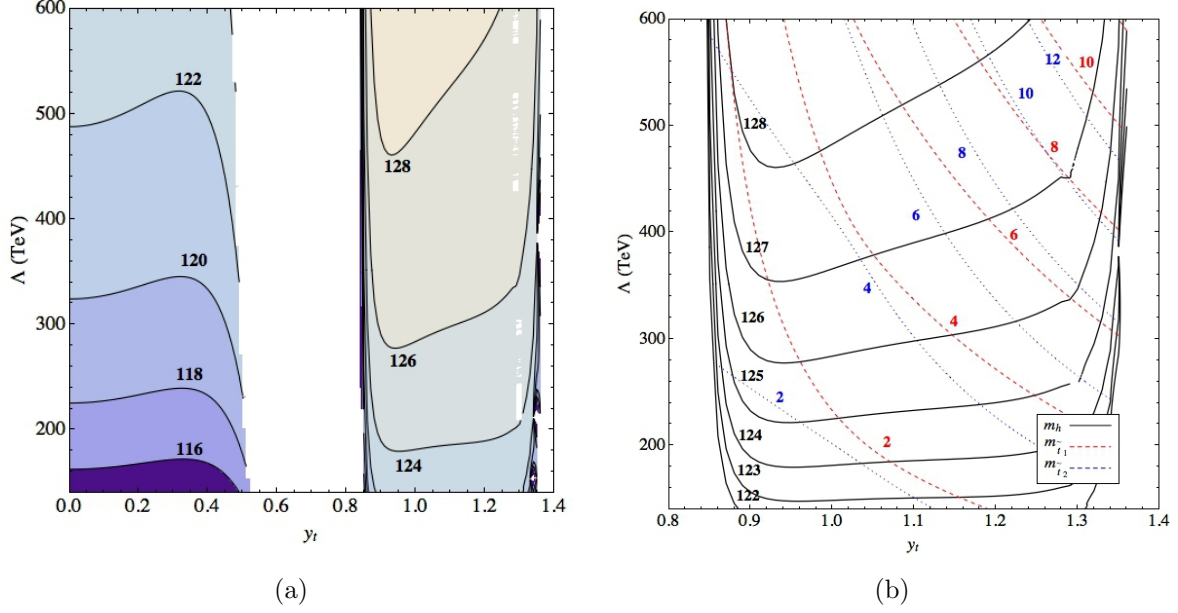


FIG. 3: Same plots as in Fig 1(a) and and 1(b) with $M = 10^8$ GeV, $\tan \beta = 20$.

at 337 GeV and the heavier stop at 1.6 TeV (the full superpartner spectrum for this choice of parameters is presented in Table II as spectrum 5).

While the models with only messenger-top Yukawas are the most economical ones, viable models with additional messenger-bottom and messenger-tau couplings are also allowed. As an example, in Fig. 4 we present Higgs mass contours in the $y_t - y_b$ plane for $N = 2$, with $M = 10^8$ GeV, $\Lambda = 230$ TeV, and $\tan \beta = 10$. For large values of y_b , the sleptons become tachyonic, leading to the white excluded region on the right.

Since we are turning on order-one superpotential couplings, it is interesting to ask at what scales these become non-perturbative. For example, if $N = 1$ the high scale models with $M = 10^{12}$ GeV remain perturbative even above the GUT scale. For $M = 10^8$ GeV,

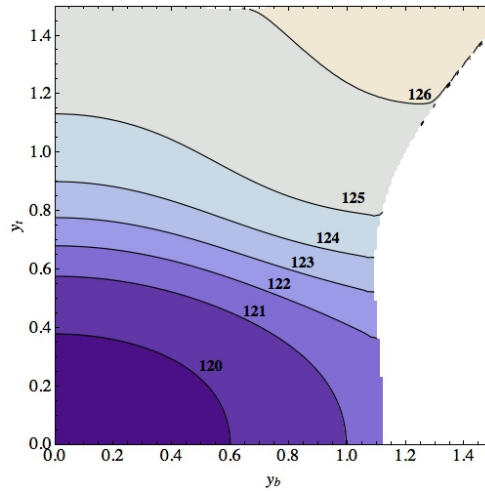


FIG. 4: Higgs mass contours in $y_t - y_b$ plane with $M = 10^8$ GeV, $\Lambda = 230$ TeV, and $\tan \beta = 10$.

parameter	spectrum 1	spectrum 2	spectrum 3	spectrum 4	spectrum 5
M_{mess}	$9 \cdot 10^5$	$9 \cdot 10^5$	$1 \cdot 10^8$	$1 \cdot 10^{12}$	$7.9 \cdot 10^{12}$
Λ	$1.29 \cdot 10^5$	$1.49 \cdot 10^5$	$2.48 \cdot 10^5$	$3.24 \cdot 10^5$	$2.40 \cdot 10^5$
$\tan \beta$	6	10	20	10	10
y_t	0.959	1.09	0.93	0.82	0.78
μ	1834	2489	3201	3187	2132
h_0	126.6	124.0	125.5	124.0	126.7
A	1929	2581	3321	3478	2376
H_0	1930	2577	3322	3479	2378
H_{\pm}	1931	2583	3322	3479	2378
t_1	348	1264	1453	746	337
t_2	1040	1606	2082	2242	1601
b_L	1017	1524	2061	2286	1647
b_R	1318	1491	2220	2632	1955
u_L, c_L	1399	1593	2431	2930	2188
u_R, c_R	1333	1501	2266	2695	2023
d_L, s_L	1401	1594	2432	2931	2189
d_R, s_R	1330	1515	2276	2665	1980
e_L, μ_L	459	500	885	1318	1027
e_R, μ_R	247	368	604	862	633
τ_L	461	511	888	1314	1024
τ_R	240	343	553	849	622
ν_e	452	493	876	1312	1021
ν_{μ}, ν_{τ}	452	494	881	1315	1024
χ_1	175	204	340	441	324
χ_2	345	402	659	847	626
χ_3	-1852	-2518	-3245	-3231	2160
χ_4	1855	2519	3249	3232	-2162
χ_1^{\pm}	345	402	659	848	626
χ_2^{\pm}	1855	2519	3250	3232	2162
g	1020	1167	1824	2298	1745

TABLE II: Model parameters, and resulting Higgs parameters and spectra for five sample models. All mass parameters are given in GeV.

one loses perturbativity at around $10^{12} - 10^{13}$ GeV for $y_t \sim 1$. Finally, with $M = 900$ TeV, the models stay perturbative up to $10^9 - 10^{10}$ GeV for y_t above 1, and for a smaller $y_t \sim 0.9$ (as is the case with the lightest spectrum we discuss below — see Table II) up to scales of around 10^{13} GeV. In the case of $N = 2$ the couplings remain perturbative for a few order of magnitude above the messenger scale.

B. Superpartner spectra and LHC signatures

To understand the phenomenology of our FGM models, we present in Table II complete

superpartner spectra for several choices of the parameters at low, intermediate and high messenger scales. A detailed analysis of the experimental signatures is beyond the scope of this paper, but it is nonetheless useful to point out a few key features.

Since the examples we studied in detail have a single pair of messengers, the NLSP is always a (bino) neutralino. Whether or not the neutralino decays inside the detector depends on the gravitino mass, which involves a lot of uncertainty⁷. In any case, the models can be probed by “standard” supersymmetry searches which are based on jets plus missing energy.

We have chosen to display here two examples (spectra 1 and 5) which contain at least one very light colored superpartner. The stop mass in spectrum 1 is 348 GeV (with a neutralino mass of 175 GeV), while the first two generation squark masses are around 1.3–1.4 TeV. Spectrum 5 has the light stop at 337 GeV, and a heavier stop at 1.6 TeV. Both of these examples are tuned in the sense that they rely on the cancellation between the pure gauge and the $y_t^2 g^2$ term in order to obtain small stop masses with a large stop mixing. Thus, small variations of the parameters may result in tachyonic stops or in the failure of EWSB. The remaining spectra of Table II do not rely on this cancellation and therefore provide more generic examples of FGM, with heavier superpartners. In all of these examples, the stop mixing is dominated by the A -term, which is larger than $\mu \cot \beta$ by roughly an order of magnitude.

It is interesting to note that we have found light spectra for both high and low messenger scales, and thus a large region of messenger scales and gravitino masses can be accommodated within viable FGM models with light superpartners. Since, with the simplifying assumption of a massless neutralino, the current limit on equal-mass squarks and gluinos is around 1.5 TeV [27], some of our models will soon be probed by LHC data.

One could also study models with more messenger pairs, which would lead to slepton NLSP’s. Since the pure gauge contribution to the scalar squared masses is proportional to N , we expect that for such models the cancellation between the $y^2 g^2$ and pure gauge contributions will be less dramatic than for the $N = 1$ models, so that very light spectra such as spectrum 1 would not occur. However, superpartner masses around 1-2 TeV, as in the remaining spectra of Table II, which do not rely on this cancellation, would be possible, with a slepton NLSP and with gluinos that are heavier than the squarks.

IV. CONCLUSIONS

We have shown that gauge-mediated models with messenger-matter couplings can give rise to an acceptable Higgs mass, with colored superpartners within the LHC reach. The important new ingredient in these models is the messenger-scale top A -term, which gives rise to significant stop mixing, and therefore enhances the Higgs mass. This mechanism is quite insensitive to the details of the basic GMSB model, and works for a wide range of messenger scales. While we have concentrated on models with only a single pair of

⁷ If F_X is the dominant F -term, the gravitino mass in our models varies between about 100 eV for $M = 900$ TeV to a GeV for the high-scale models. It is quite plausible however that there are much higher F terms in the supersymmetry-breaking sector, as is often the case in calculable models, in which case F_X is generated through several small couplings from a higher F -term. For large values of the gravitino mass, the neutralino would decay outside the detector. Heavy gravitinos from late NLSP decay could also supply warm dark matter [7, 26].

messengers and a messenger-top coupling, the results can be easily generalized to models with several messenger pairs, with or without messenger down-type couplings. The resulting spectra can therefore be quite diverse, with either a neutralino or a slepton NLSP. We leave a detailed study of the phenomenology of these models for future work. It would be particularly interesting to examine models with down-type messenger couplings, which may generate sufficient stau mixing and thus enhance the Higgs to $\gamma\gamma$ branching ratio [28].

The structure of the matter-messenger couplings can naturally be the same as the structure of the usual Yukawas, since the new couplings are obtained by replacing H_U and/or H_D by the messenger of the same gauge charge. The models are therefore protected against flavor-violation by supersymmetric alignment. This can be simply realized in the context of flavor symmetries if the Higgs and corresponding messenger have the same charges under the flavor symmetry.

The amount of tuning in our models is related to the tuning of the new messenger-matter coupling y_t . As can be seen in Table II, for spectrum 1, which relies on a finely-dialed y_t to obtain a light spectrum, the μ -term is 1.8 TeV, while less tuned choices of y_t require larger μ -terms, around 2 or 3 TeV. It is interesting to see the effect of the 126 GeV Higgs mass on the amount of tuning required, as compared with a pure GMSB model with a lower Higgs mass. Thus for example, in spectrum 3 of Table II, the μ -term is 3.2 TeV, whereas the same mGMSB model with $y_t = 0$, which gives a Higgs mass of 119 GeV, has $\mu = 1.1$ TeV. Thus, the Higgs mass being 126 GeV and not 119 GeV worsens the amount of fine tuning by roughly a factor of 8. Generating an acceptable μ -term in our models is an important direction to pursue. If there is a successful mechanism for generating the μ -term, the tuning question would translate into the question of how finely one needs to tune the coupling y_t . It is probably far from trivial to find a successful mechanism for generating the μ -term, and even if it is found, the required tuning of y_t is likely to be significant. Still, since y_t is a superpotential coupling, the tuning involved is qualitatively different from the tuning of the Higgs mass in the standard model.

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Appendix A: Derivation of the soft terms

In this Appendix, we outline the calculation of the soft terms from wave-function renormalizations. Our discussion closely follows [19, 20, 29], but we pay special attention to two issues: the first is the one-loop mixing between messengers and SM fields, and the second is the matrix structure of the couplings and soft terms in generation space.

1. Messenger-SM Mixing

The 2-loop soft terms can be extracted, to leading order in F/M^2 , from wave function renormalizations [19]. The main advantage of this method is that the 2-loop expressions are given by combinations of the beta functions and anomalous dimensions of the supersymmetric theory at one loop.

Superpotential couplings of the messengers to the SM fields complicate the calculation in two ways. First, in general one cannot integrate the RGE equations to obtain analytic expressions for the wave-function renormalizations and anomalous dimensions at different scales. Second, starting at one-loop, the messengers and SM fields mix.

One can capture all the relevant ingredients of the calculation in a single-generation toy model with messenger D couplings to leptons, with the superpotential,

$$W = XD\bar{D} + yDle + YHle. \quad (\text{A1})$$

At low scales, this model has two types of fields: l and e , which couple to the messenger D but cannot mix with it, and H , which has no coupling to D but can mix with it. We will sketch the computation of the soft mass for each.

The soft masses of any field $\phi = H, l, e$, can be extracted from its wave-function renormalization $Z_\phi(\mu, M)$ at $\mu \leq M$, by promoting the threshold M to $X = M + F\theta^2$. Just above the messenger scale M , the Kähler potential for D, H , is of the form,

$$K = D^\dagger D(1 + \dots) + H^\dagger H(1 + \dots) + a_{DH} D^\dagger H + \text{h.c.} \quad (\text{A2})$$

where the mixing arises from the running between the cutoff scale Λ and M , $a_{DH} \propto \gamma_{DH} \log(M/\Lambda)$, and $\gamma_{DH} \sim 1/(16\pi^2)Y^\dagger y$. Clearly, however, this mixing cannot affect the derivation of the soft terms, since these only result from integrating out the messengers at the threshold X . Note also that we are interested in the soft masses at 2-loops, so the only relevant part of the mixing is the one-loop contribution. This is certainly supersymmetric, since the fields running in the loop are just l and e .

Imagine then that we define new combinations, \tilde{H} and \tilde{D} , of H and D such that, just above the messenger scale, the Kähler potential is canonical in terms of \tilde{H}, \tilde{D} . Both these fields will couple to X in the superpotential, but only a single combination of them (this combination has a small admixture of H), gets mass. We can then rotate to the mass basis, so that the superpotential is

$$W = X\tilde{D}\bar{\tilde{D}} + \tilde{y}\tilde{D}le + \tilde{Y}\tilde{H}le, \quad (\text{A3})$$

and the Kähler potential is still canonical in terms of \tilde{D} and \tilde{H} . The new couplings \tilde{Y} and \tilde{y} differ from the original couplings, but the difference is one-loop order and therefore will not affect the results.

Let us consider first the soft mass of l . We have

$$\ln Z_l(t) = \int_{t_0}^M dt' \gamma_l^>(t') + \int_M^t dt' \gamma_l^<(t') \quad (\text{A4})$$

so

$$\frac{\partial \ln Z_l(t)}{\partial \ln |X|} = \gamma_l^>(\ln M) - \gamma_l^<(\ln M) + \int_{|X|}^t dt' \frac{\partial \gamma_l^<(t')}{\partial \ln |X|} \quad (\text{A5})$$

because the anomalous dimension has a jump at M . Then

$$\frac{\partial^2 \ln Z_l(t)}{\partial \ln |X|^2} = \sum_{\lambda} \frac{\partial \Delta \gamma_l^>}{\partial \lambda}(M) \beta_{\lambda}(M) - \frac{\partial \gamma_l^<(t')}{\partial \ln |X|} \Big|_{X=M} \quad (\text{A6})$$

where λ sums over all couplings—gauge, y , and Y , and where we used the fact that another derivative would make the expression 3-loops since the integrand is already 2-loop. Finally we can write this as

$$\frac{\partial^2 \ln Z_l(t)}{\partial \ln |X|^2} = \sum_{\lambda} \frac{\partial \Delta \gamma_l^>}{\partial \lambda}(M) \beta_{\lambda}(M) - \sum_{\lambda} \frac{\partial \gamma_l^<(t')}{\partial \lambda} \Delta \beta_{\lambda} \Big|_{X=M} \quad (\text{A7})$$

Note that y will only appear in the first term.

Now consider H instead. Its anomalous dimension is continuous at the threshold so we just have,

$$\frac{\partial \ln Z_H(t)}{\partial \ln |X|} = \int_{|X|}^t dt' \frac{\partial \gamma_H^<(t')}{\partial \ln |X|} \quad (\text{A8})$$

and therefore

$$\frac{\partial^2 \ln Z_H(t)}{\partial \ln |X|^2} = - \sum_{\lambda} \frac{\partial \gamma_H^<(t')}{\partial \lambda} \Delta \beta_{\lambda} \Big|_{X=M} \quad (\text{A9})$$

Let us now return to the $D - H$ mixing. The beta functions which appear in the final expressions (A7), (A9) contain various anomalous dimensions. In particular, before the field redefinition discussed above, they contain the anomalous dimension γ_{DH} . As we saw above, however, to obtain the soft terms, one can consistently work with canonically normalized fields. The “mixed” anomalous dimension $\gamma_{DH}^>$ should therefore be set to zero⁸. While some mixing is generated above the messenger scale, its effects on the soft terms are of higher order in the loop expansion

2. Final expressions

We are now ready to present the soft terms resulting from general coupling matrices y_U , y_D and y_L .

We will first discuss the 2-loop results, and then turn to the 1-loop F^4/M^6 contributions. As outlined above, at leading order in F/M^2 , the soft terms are given by,

$$\tilde{m}^2 = -\frac{1}{4} \sum_{\lambda} \left(\frac{d\Delta\gamma}{d\lambda} \beta_{>} - \frac{d\gamma_{<}}{d\lambda} \Delta\beta \right) \left| \frac{F}{M} \right|^2, \quad (\text{A10})$$

where $\Delta\gamma \equiv \gamma_{>} - \gamma_{<}$ and $\Delta\beta \equiv \beta_{>} - \beta_{<}$, and

$$A_{abc} = \frac{1}{2} \left(\lambda_{a'bc} \Delta\gamma_{a'}^{a'} + \lambda_{ab'c} \Delta\gamma_b^{b'} + \lambda_{abc'} \Delta\gamma_c^{c'} \right) \frac{F}{M}. \quad (\text{A11})$$

⁸ Our results for the Higgs soft masses differ from [20] due to this effect.

Taking into account the full generation structure the A-terms at $\mu = M$ take the form

$$\begin{aligned} A_U &= -\frac{1}{16\pi^2} \left[\left(y_U y_U^\dagger + y_D y_D^\dagger \right) Y_U + 2Y_U \left(y_U^\dagger y_U \right) \right] \frac{F}{M} \\ A_D &= -\frac{1}{16\pi^2} \left[\left(y_U y_U^\dagger + y_D y_D^\dagger \right) Y_D + 2Y_D \left(y_D^\dagger y_D \right) \right] \frac{F}{M} \\ A_L &= -\frac{1}{16\pi^2} \left[\left(y_L y_L^\dagger \right) Y_L + 2Y_L \left(y_L^\dagger y_L \right) \right] \frac{F}{M}, \end{aligned} \quad (\text{A12})$$

and the 2-loop soft squared masses at $\mu = M$ are

$$\tilde{m}_{H^u}^2 = \frac{|F/M|^2}{(4\pi)^4} \left\{ -3Tr \left(Y_U^\dagger y_U y_U^\dagger Y_U + Y_U^\dagger y_D y_D^\dagger Y_U + 2Y_U^\dagger Y_U y_U^\dagger y_U \right) + 2N \left(\frac{3}{4}g_2^4 + \frac{3}{20}g_1^4 \right) \right\} \quad (\text{A13})$$

$$\begin{aligned} \tilde{m}_{H^d}^2 &= \frac{|F/M|^2}{(4\pi)^4} \left\{ -3Tr \left(Y_D^\dagger y_U y_U^\dagger Y_D + Y_D^\dagger y_D y_D^\dagger Y_D + 2Y_D^\dagger Y_D y_D^\dagger y_D \right) \right. \\ &\quad \left. - Tr \left(Y_L^\dagger y_L y_L^\dagger Y_L + 2Y_L^\dagger Y_L y_L^\dagger y_L \right) + 2N \left(\frac{3}{4}g_2^4 + \frac{3}{20}g_1^4 \right) \right\} \end{aligned} \quad (\text{A14})$$

$$\begin{aligned} \tilde{m}_q^2 &= \frac{|F/M|^2}{(4\pi)^4} \left\{ \left(3Tr \left(y_U^T y_U^* \right) - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{13}{15}g_1^2 \right) y_U^* y_U^T \right. \\ &\quad \left(Tr \left(3y_D^T y_D^* + y_L^T y_L^* \right) - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{7}{15}g_1^2 \right) y_D^* y_D^T \\ &\quad + 3y_U^* y_U^T y_U^* y_U^T + 3y_D^* y_D^T y_D^* y_D^T + y_D^* y_D^T y_U^* y_U^T + y_U^* y_U^T y_D^* y_D^T \\ &\quad + 2y_U^* Y_U^T Y_U^* y_U^T - 2Y_U^* y_U^T y_U^* Y_U^T + 2y_D^* Y_D^T Y_D^* y_D^T - 2Y_D^* y_D^T y_D^* Y_D^T \\ &\quad \left. + 2N \left(\frac{4}{3}g_3^4 + \frac{3}{4}g_2^4 + \frac{1}{60}g_1^4 \right) 1_{3 \times 3} \right\} \end{aligned} \quad (\text{A15})$$

$$\begin{aligned} \tilde{m}_{u^c}^2 &= \frac{|F/M|^2}{(4\pi)^4} \left\{ 2 \left(3Tr \left(y_U^\dagger y_U \right) - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{13}{15}g_1^2 \right) y_U^\dagger y_U \right. \\ &\quad + 6y_U^\dagger y_U y_U^\dagger y_U + 2y_U^\dagger y_D y_D^\dagger y_U + 2y_U^\dagger Y_U Y_U^\dagger y_U - 2Y_U^\dagger y_U y_U^\dagger Y_U \\ &\quad \left. + 2y_U^\dagger Y_D Y_D^\dagger y_U - 2Y_U^\dagger y_D y_D^\dagger Y_U + 2N \left(\frac{4}{3}g_3^4 + \frac{4}{15}g_1^4 \right) 1_{3 \times 3} \right\} \end{aligned} \quad (\text{A16})$$

$$\begin{aligned}\tilde{m}_{dc}^2 = \frac{|F/M|^2}{(4\pi)^4} \Big\{ & 2 \left(\text{Tr} \left(3y_D^\dagger y_D + y_L^\dagger y_L \right) - \frac{16}{3}g_3^2 - 3g_2^2 - \frac{7}{15}g_1^2 \right) y_D^\dagger y_D \\ & + 6y_D^\dagger y_D y_D^\dagger y_D + 2y_D^\dagger y_U y_U^\dagger y_D + 2y_D^\dagger Y_U Y_U^\dagger y_D + 2y_D^\dagger Y_D Y_D^\dagger y_D \\ & - 2Y_D^\dagger y_U y_U^\dagger Y_D - 2Y_D^\dagger y_D y_D^\dagger Y_D + 2N \left(\frac{4}{3}g_3^4 + \frac{1}{15}g_1^4 \right) 1_{3 \times 3} \Big\}\end{aligned}\quad (\text{A17})$$

$$\begin{aligned}\tilde{m}_L^2 = \frac{|F/M|^2}{(4\pi)^4} \Big\{ & \left(\text{Tr} \left(3y_D^T y_D^* + y_L^T y_L^* \right) - 3g_2^2 - \frac{9}{5}g_1^2 \right) y_L^* y_L^T \\ & + 3y_L^* y_L^T y_L^* y_L^T + 2y_L^* Y_L^T Y_L^* y_L^T - 2Y_L^* y_L^T y_L^* Y_L^T \\ & + 2N \left(\frac{3}{4}g_2^4 + \frac{3}{20}g_1^4 \right) 1_{3 \times 3} \Big\}\end{aligned}\quad (\text{A18})$$

$$\begin{aligned}\tilde{m}_{ec}^2 = \frac{|F/M|^2}{(4\pi)^4} \Big\{ & 2 \left(\text{Tr} \left(3y_D^\dagger y_D + y_L^\dagger y_L \right) - 3g_2^2 - \frac{9}{5}g_1^2 \right) y_L^\dagger y_L \\ & + 6y_L^\dagger y_L y_L^\dagger y_L + 2y_L^\dagger Y_L Y_L^\dagger y_L - 2Y_L^\dagger y_L y_L^\dagger Y_L + \frac{6}{5}N g_1^4 1_{3 \times 3} \Big\}.\end{aligned}\quad (\text{A19})$$

In addition, we give here the fully-flavored 1-loop contributions to the soft masses,

$$\delta m_q^2 = -\frac{1}{(4\pi)^2} \frac{1}{6} \left(y_U^* y_U^T + y_D^* y_D^T \right) \frac{F^4}{M^6} \quad (\text{A20})$$

$$\delta m_{uc}^2 = -\frac{1}{(4\pi)^2} \frac{1}{3} \left(y_U^\dagger y_U \right) \frac{F^4}{M^6} \quad (\text{A21})$$

$$\delta m_{dc}^2 = -\frac{1}{(4\pi)^2} \frac{1}{3} \left(y_D^\dagger y_D \right) \frac{F^4}{M^6} \quad (\text{A22})$$

$$\delta m_l^2 = -\frac{1}{(4\pi)^2} \frac{1}{6} \left(y_L^* y_L^T \right) \frac{F^4}{M^6} \quad (\text{A23})$$

$$\delta m_{ec}^2 = -\frac{1}{(4\pi)^2} \frac{1}{3} \left(y_L^\dagger y_L \right) \frac{F^4}{M^6}. \quad (\text{A24})$$

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